# A Constituent Picture of Hadrons from Light-Front QCD

**Robert J. Perry**
*Department of Physics, The Ohio State University, Columbus, Ohio, USA*

It may be possible to derive a constituent approximation for bound states in QCD using hamiltonian light-front field theory. Cutoffs that violate explicit gauge invariance and Lorentz covariance must be employed. A similarity renormalization group and coupling coherence are used to compute the effective hamiltonian as an expansion in powers of the canonical QCD running coupling constant. At second order the QCD hamiltonian contains a confining interaction, which is being studied using bound state perturbation theory. Explicit constituent masses appear because of symmetry violations, and confinement also produces mass gaps, leading to the possibility of an accurate non-perturbative constituent approximation emerging in light-front QCD.

## 1. Introduction

The solution of Quantum Chromodynamics in the non-perturbative domain remains one of the most important and interesting unsolved problems in physics. QCD is believed to be the fundamental theory of the strong interaction, but even its definition in the non-perturbative domain is problematic. There are many sources of difficulty, but they can all be traced to the fact that QCD is formulated as a theory of an infinite number of degrees of freedom that span an infinite number of energy scales.

The basic assumption upon which all of our work is based is that it is possible to *derive a constituent picture for hadrons from QCD* [1-5]. If this is possible, non-perturbative bound state problems in QCD are approximated as coupled, few-body Schrödinger equations. For a meson, we then have,

$$P^- \mid \Psi\rangle = \frac{P_\perp^2 + M^2}{P^+} \mid \Psi\rangle, \tag{1}$$

where,

$$\mid \Psi\rangle = \phi_{q\bar{q}} \mid q\bar{q}\rangle + \phi_{q\bar{q}g} \mid q\bar{q}g\rangle + \cdots. \tag{2}$$

$P^-$ is the light-front hamiltonian, $P_\perp$ is the total transverse momentum, $P^+$ is the total longitudinal momentum, and $M$ is the invariant mass of the state. We assume that to 'leading order' a low-lying meson can be approximated as a quark/antiquark pair, with additional quarks and gluons producing 'perturbative' corrections that can be systematically computed.

Many severe problems must be overcome to arrive at this formulation of the bound state problem; however, the final advantages are huge. The result is a formulation of the non-perturbative problem in a form directly accessible to physical intuition, which has proven essential for guiding approximations in atomic calculations. Variational methods

and large matrix diagonalization are powerful numerical tools that can be used after the hamiltonian is determined.

I must emphasize that *it is not our intent to simply force the constituent approximation on the theory by employing a Tamm-Dancoff truncation on the number of particles ab initio.* We worked on such an approach initially [6], and gained valuable insights; but it became clear that we have no good method of controlling the nonlocalities resulting from particle number truncation without a dynamical mechanism that naturally limits the number of particles in a state.

Any student of field theory should immediately be suspicious of the possibility that a constituent approximation can arise, although QED provides an important accepted example that guides much of our work [2]. How can a constituent approximation arise in any field theory?

Fock space is extremely large, an infinite sum of cross products of infinite dimensional Hilbert spaces. It is not obvious that the low-lying eigenstates should have significant support only in the few-body sectors of Fock space. In fact, this simply does not happen in perturbation theory. In perturbation theory high-energy many-body states do not decouple from low-energy few-body states. Consider an electron mixing with high-energy electron/photon states. The error made by simply throwing away the high energy components of the state is infinite. Moreover, there are an infinite number of scales and both the electron and photon that 'dress' the low-energy bare electron are in turn dressed by additional pairs, *ad infinitum.*

The lesson here is quite old. Without *regularization and renormalization* a constituent picture is impossible. Renormalization may allow us to move the dynamical effects of high-energy, many-body states from the eigenstate to effective interactions between effective quarks and gluons.

Low-energy many-body states also do not decouple from low-energy few-body states. In fact, it is common lore that hadrons are excitations on an extremely complicated vacuum. Students of QCD expect the infinite-body vacuum to be an integral part of every hadron eigenstate. This is the problem that leads us to use light-front coordinates, just as it motivated the use of the infinite momentum frame for the formulation of the parton model. In light-front coordinates physical particle trajectories satisfy the kinematic relativistic constraint

$$p^+ \geq 0 \; , \tag{3}$$

because all velocities are equal to or less than the velocity of light. Since longitudinal momentum is conserved, the only states that can mix with the zero momentum bare vacuum are those in which every bare parton has identically zero longitudinal momentum. For a free particle of mass $m$, the light-front energy is

$$p^- = \frac{\mathbf{p}_\perp^2 + m^2}{p^+} \; . \tag{4}$$

This energy diverges as $p^+$ approaches zero, which must happen as the number of particles grows for fixed total longitudinal momentum. Thus, in light-front coordinates all many-body states become high energy states, leading us back to the original problem

of replacing the effects of high energy states with effective interactions. This argument is naive, but there is little profit in elaborating further at this point.

Finally, manifest gauge invariance and manifest covariance apparently require all states to contain an infinite number of particles. This is most easily seen, for example, by considering rotation operators. Rotations are dynamical in light-front coordinates and the generators contain interactions that change particle number. No state with a finite number of particles transforms correctly under rotations. We use *cutoffs that violate these symmetries*, which must then be repaired by effective interactions that remove all cutoff dependence. The constituent approximation is possible only if these symmetries are also treated approximately. Proposing the violation of manifest gauge invariance is heresy in the QCD community, but heresy sometimes leads to progress in science.

There is a long list of questions concerning how a constituent approximation can arise in QCD, but I mention only one. How can confinement emerge without a complicated vacuum? Since we have a hamiltonian, we can use a variational calculation to study what happens as a quark/antiquark pair are separated to infinity. Since the addition of gluons can only lower the energy, we must find that

$$\langle \phi_{q\overline{q}} \mid H \mid \phi_{q\overline{q}} \rangle \longrightarrow \infty \;\; \text{as} \;\; R \rightarrow \infty \; . \tag{5}$$

Here $R$ is the quark separation, and the only way this matrix element can diverge is if the hamiltonian contains a two-body interaction that diverges. We will see below that this apparently happens.

This discussion is not intended to convince the skeptical reader that a constituent approximation is valid. However, *the assumption that a constituent picture emerges from QCD provides strong guidance.* A hamiltonian approach is indicated. Cutoffs that limit mixing of high and low energy states are required, and they must violate explicit rotational covariance and gauge invariance. All non-perturbative effects attributed to the vacuum in other approaches must directly appear in few-body effective interactions.

## 2. Light-Front Renormalization Group

The renormalization of the hamiltonian and all other dynamical observables begins with the observation that no physical result can depend on the cutoff. In the Schrödinger equation,

$$P^-_\Lambda \mid \Psi_\Lambda \rangle = \frac{P_\perp^2 + M^2}{P^+} \mid \Psi_\Lambda \rangle \; , \tag{6}$$

the eigenvalue, $M$, cannot depend on the cutoff. The hamiltonian, $P^-$, must depend on the cutoff, as does the eigenstate. Wilson's renormalization group was formulated starting with the observation that physical matrix elements cannot depend on the cutoff, and we have adapted his approach to the light-front problems we face [7].

It is not possible to discuss cutoff-independence if the cutoff is fixed, so the central operator in Wilson's renormalization group is a transformation that lowers the cutoff. Given a transformation, $T$, that lowers the cutoff by a factor of 1/2, for example, we can

define a renormalized hamiltonian to be one which has a finite cutoff but results from an infinite number of transformations. The transformation determines what operators must be precisely controlled for this limit to exist. Near a fixed point (*i.e.*, a hamiltonian that does not change under the transformation), these operators can be classified as relevant and marginal.

In the perturbative regime relevant and marginal operators are determined by their naive engineering dimension. In light-front field theory there is no longitudinal locality, only transverse locality, so it is the transverse dimension of an operator that determines its classification. However, while there are a finite number of relevant and marginal operators in equal-time field theory, the violation of longitudinal locality in light-front field theory implies that ratios of longitudinal momenta can appear, allowing entire functions of longitudinal momentum fractions to appear in each relevant and marginal operator. At first sight this appears to be a disaster; however, one paradox we faced above was how complicated interactions associated with non-perturbative effects such as confinement could arise in few-body operators. This is possible because of the violation of longitudinal locality.

To develop a light-front renormalization group we must decide what cutoff to implement and then derive a transformation that runs this cutoff. It is possible to use a cutoff on the total invariant-mass of states, as is commonly done in DLCQ calculations for example; however, such cutoffs lead to strong spectator dependence and small energy denominators appear in the resultant effective interactions. We use a cutoff on the change in free energy. If the hamiltonian is viewed as a matrix such a cutoff limits how far off diagonal matrix elements can appear.

There is not enough space to elaborate the transformation that runs this cutoff, so I refer the reader to the literature [1,2,4,8,9]. The transformation is unitary, leading to what Głazek and Wilson call a *similarity renormalization group* [8,9]. Let $H = h_0 + v$, where $h_0$ is a free hamiltonian with $h_0 \mid \phi_i\rangle = E_{0i} \mid \phi_i\rangle$, and $v$ is cut off so that

$$\langle\phi_i|v|\phi_j\rangle = 0 \;, \tag{7}$$

if $|E_{0i} - E_{0j}| > \Lambda$. If this cutoff is lowered to $\Lambda'$, the new hamiltonian matrix elements to $\mathcal{O}(v^2)$ are

$$\begin{aligned} H'_{ab} \quad = \quad & \langle\phi_a|h_0+v|\phi_b\rangle - \sum_k v_{ak} v_{kb} \Bigg[ \frac{\theta\big(|\Delta_{ak}|-\Lambda'\big)\theta\big(|\Delta_{ak}|-|\Delta_{bk}|\big)}{E_{0k}-E_{0a}} \\ & \qquad\qquad + \frac{\theta\big(|\Delta_{bk}|-\Lambda'\big)\theta\big(|\Delta_{bk}|-|\Delta_{ak}|\big)}{E_{0k}-E_{0b}} \Bigg], \end{aligned} \tag{8}$$

where $\Delta_{ij} = E_{0i} - E_{0j}$ and $|E_{0a} - E_{0b}| < \Lambda'$. To follow the details of the discussion it is important to remember that there are implicit cutoffs in this expression because the matrix elements of $v$ have already been cut off so that $v_{ij} = 0$ if $|E_{0i} - E_{0j}| > \Lambda$.

It is rather easy to understand this result qualitatively. We have removed the coupling between degrees of freedom whose free energy difference is between $\Lambda'$ and $\Lambda$,

so the effects of these couplings are forced to appear in the new hamiltonian as direct interactions. To first order, the new hamiltonian is the same as the old hamiltonian, except that couplings of states with energy differences between $\Lambda'$ and $\Lambda$ are now zero. To second order, the new hamiltonian contains a new interaction which sums over the second-order effects of couplings that have been removed. The second-order term in the new hamiltonian resembles the expression found in second-order perturbation theory, which is not surprising since the new hamiltonian must produce the same perturbative expansion for eigenvalues, cross sections, etc. as the original hamiltonian.

Equation (9) shows how the hamiltonian changes when the cutoff is lowered, and the next step is to determine from this change what hamiltonians can emerge from an infinite number of transformations. The simplest result is a *fixed point* hamiltonian, one which does not change under the transformation. In $3+1$ dimensions the only known fixed points are free field theories. *Coupling coherence* is a generalization of the fixed point idea [2,7,10]. A coupling coherent hamiltonian reproduces itself in form, but one or more couplings run while all additional couplings are invariant functions of these running couplings. For example, in QCD the canonical coupling runs at third order. To second order in this coupling, all interactions must reproduce themselves exactly, with $\Lambda \rightarrow \Lambda'$. It is not trivial to implement this simple-sounding constraint, but at each order it determines the hamiltonian. In all calculations to date the resultant hamiltonian is unique, and all broken symmetries are restored to the order at which the hamiltonian is fixed.

To second order a generic coupling coherent hamiltonian that contains $v$ must also contain,

$$
\begin{aligned}
H_{ab} \;=\; \langle\phi_a|h_0+v|\phi_b\rangle - \sum_k v_{ak}v_{kb}\Bigg[&\frac{\theta\big(|\Delta_{ak}|-\Lambda\big)\theta\big(|\Delta_{ak}|-|\Delta_{bk}|\big)}{E_{0k}-E_{0a}} \\
&+\frac{\theta\big(|\Delta_{bk}|-\Lambda\big)\theta\big(|\Delta_{bk}|-|\Delta_{ak}|\big)}{E_{0k}-E_{0b}}\Bigg], \qquad (9)
\end{aligned}
$$

or

$$
\begin{aligned}
H_{ab} \;=\; \langle\phi_a|h_0+v|\phi_b\rangle + \sum_k v_{ak}v_{kb}\Bigg[&\frac{\theta\big(\Lambda-|\Delta_{ak}|\big)\theta\big(|\Delta_{ak}|-|\Delta_{bk}|\big)}{E_{0k}-E_{0a}} \\
&+\frac{\theta\big(\Lambda-|\Delta_{bk}|\big)\theta\big(|\Delta_{bk}|-|\Delta_{ak}|\big)}{E_{0k}-E_{0b}}\Bigg]. \qquad (10)
\end{aligned}
$$

Note that $v$ in these expressions is the same as that above only to first order. The coupling coherent interaction in $H$ is written as a power series in $v$ which reproduces itself under the transformation, except the cutoff changes. In higher orders the canonical variables also run.

The light-front similarity renormalization group and coupling coherence fix the QCD hamiltonian as an expansion in powers of the running canonical coupling.

## 3. QCD: A Strategy for Bound State Calculations and Confinement

While realistic calculations will no doubt require a more elaborate procedure, a relatively simple strategy for doing bound state calculations can now be outlined [2,4].

> i)Start with the canonical hamiltonian, $H_{can}$, and use the similarity renormalization group and coupling coherence to compute,
>
> $$H^\Lambda = h_0^\Lambda + g_\Lambda h_1^\Lambda + g_\Lambda^2 h_2^\Lambda + \cdots \tag{11}$$
>
> Truncate this series at a fixed order.
>
> ii) Choose an approximate hamiltonian that can be treated non-perturbatively,
>
> $$H^\Lambda = H_0^\Lambda + V^\Lambda \; . \tag{12}$$
>
> You must choose $\Lambda$ and $H_0^\Lambda$ to minimize errors.
>
> iii) Accurately solve $H_0^\Lambda$ as the leading approximation.
>
> iv) Compute higher order corrections from $V^\Lambda$ using bound state perturbation theory.
>
> v) To improve the calculation further return to step (i) and compute the hamiltonian to higher order.

There are two principal reasons that this strategy will fail for QCD. First, the hamiltonian is computed perturbatively so that errors in the strengths of all operators are at least as large as a power of $\alpha$. Small errors in the strengths of irrelevant operators tend to produce even smaller errors in results. However, errors in marginal operators tend to produce errors of the same order in results and small errors in relevant operators tend to produce exponentially large errors in results. At the minimum we expect that we will have to fine tune relevant operators, which means tuning a finite number of functions of longitudinal momenta. Second, chiral symmetry breaking operators (where light-front chiral symmetry should be distinguished from equal-time chiral symmetry [1]) will not arise at any order in an expansion in powers of the strong coupling constant. We must work in the broken symmetry phase of QCD *ab initio* and insert chiral symmetry breaking operators. Simple arguments lead us to expect that only relevant operators need to be considered if transverse locality is maintained, but there are no strong arguments for transverse locality in these operators.

Despite these limitations, this strategy may be applied to the study of bound states containing at least one heavy quark [5], as discussed by Martina Brisudová in these proceedings; although even here masses should be tuned, as expected. The strategy is conceptually simple and there are no *ad hoc* assumptions.

The first step is to compute the effective QCD hamiltonian to order $\alpha$. I refer the reader to the literature for details on the canonical hamiltonian [11]. The first applications of the approach are to mesons [5], and we assume that for sufficiently small cutoffs

we can choose $H_0^\Lambda$ to contain only interactions in $H^\Lambda$ that do not involve particle production or annihilation, as dictated by our initial assumption that a constituent picture will arise. This means we can first focus on operators that act in the quark/antiquark sector. I emphasize that all operators must be computed, and without the confining interactions in sectors containing gluons the entire approach would make no sense.

First consider the second-order correction to the quark self-energy. This results from the quark mixing with quark-gluon states whose energy is above the cutoff. If we assume that the light-front energy transfer through the quark-gluon vertex must be less than $\Lambda^2/\mathcal{P}^+$, the coupling coherent self-energy for quarks with zero current mass is

$$\Sigma_\Lambda(p) \quad = \quad \frac{g_\Lambda^2 C_F \Lambda^2}{4\pi^2 \mathcal{P}^+} \left\{ \ln\left( \frac{p^+}{\epsilon \mathcal{P}^+} \right) - \frac{3}{4} \right\} + \mathcal{O}(\epsilon) \; . \tag{13}$$

Here the quark has longitudinal momentum $p^+$, while the longitudinal momentum scale in the cutoff is $\mathcal{P}^+$. The first and most interesting feature of this result is that I have been forced to introduce a second cutoff,

$$p_i^+ > \epsilon \mathcal{P}^+ \; , \tag{14}$$

which restricts how small the longitudinal momenta of any particle can become. Without this second cutoff on the loop momenta, the self-energy is infinite, even with the vertex cutoff. This second cutoff should be thought of as a longitudinal resolution. As we let $\epsilon \rightarrow 0$ we resolve more and more wee partons, and in the process we should confront effects normally ascribed to the vacuum. In this case the wee gluons are responsible for giving the quark a mass that is literally infinite. Theorists who insist on deriving intuition from manifestly gauge invariant calculations may find this interpretation repugnant, but within the framework of a light-front hamiltonian calculation it is quite natural.

This second, infrared cutoff poses a problem. If we introduce a second cutoff, shouldn't we introduce a second renormalization group transformation to run this cutoff and find the new counterterms required by it? No. *I will insist that all divergences associated with $\epsilon \rightarrow 0$ cancel exactly in all physical results for color singlet states.* The important question is how can these divergences cancel so that mesons have a finite mass, and the answer to this question leads to confinement.

A nearly identical calculation leads to the second-order self-energy of gluons, and the dominant term goes like

$$\frac{g_\Lambda^2 \Lambda^2}{\mathcal{P}^+} \ln\left( \frac{p^+}{\epsilon \mathcal{P}^+} \right) \; . \tag{15}$$

The quark and gluon masses are infinite, which is half of the confinement mechanism. In addition to one-body operators we find quark-quark, quark-gluon, and gluon-gluon interactions. As we lower the cutoff, we remove gluon exchange interactions, and these are replaced by direct interactions. The analysis of all of these interactions is nearly

identical, and I consider only the quark-antiquark interaction. This interaction includes two pieces, instantaneous gluon exchange which is in the canonical hamiltonian, and an effective interaction resulting from high energy gluon exchange. To study confinement we need to examine the longest range part of the total interaction, which is a piece that diverges as longitudinal momentum exchange goes to zero. I outline the calculation [2,4].

High energy gluon exchange cancels part of the instantaneous gluon exchange interaction, leaving

$$V_{singular} = -4g_\Lambda^2 C_F \sqrt{p_1^+ p_2^+ k_1^+ k_2^+} \left(\frac{1}{q^+}\right)^2 \theta\big(\Lambda^2/\mathcal{P}^+ - |q^-|\big)\, \theta\big(|q^+| - \epsilon \mathcal{P}^+\big). \tag{16}$$

Here the initial and final quark (antiquark) momenta are $p_1$ and $p_2$ ($k_1$ and $k_2$), and the exchanged gluon momentum is $q$. The energies are all determined by the momenta, $p_1^- = p_{\perp 1}^2/p_1^+$, etc. This part of the interaction is independent of the spins. If $\Lambda \approx \Lambda_{QCD}$, we expect further gluon exchange to be suppressed, and we are left with this singular interaction between the quark and antiquark.

The next step in the analysis is to take the expectation value of this interaction between arbitrary quark-antiquark states, $\langle \Psi_2 \mid V \mid \Psi_1 \rangle$. If we define

$$Q = \frac{p_1 + p_2}{2} \;\; , \;\; q = p_1 - p_2 \; , \tag{17}$$

and expand the wave functions about $q = 0$, we find a divergence in the expectation value,

$$\langle \Psi_2 | V_{singular} | \Psi_1 \rangle = -\frac{g_\Lambda^2 C_F \Lambda^2}{2\pi^2 \mathcal{P}^+} \log\Big(\frac{1}{\epsilon}\Big) \int \frac{dQ^+ d^2 Q_\perp}{16\pi^3} \phi_2^*(Q) \phi_1(Q). \tag{18}$$

Unless $\phi_1$ and $\phi_2$ are the same, this vanishes by orthogonality. If they are the same, this is exactly the same expression we obtain for the expectation value of the quark plus antiquark divergent mass operators; except with the opposite sign. Therefore, there is a divergence in the quark-antiquark interaction that is independent of their relative motion and which exactly cancels the divergent masses! These cancellations only occur for color singlets, and they occur for any color singlet state with an arbitrary number of quarks and gluons. Moreover, these cancellations appear directly in the hamiltonian matrix elements, so we can take the $\epsilon \to 0$ limit before diagonalizing the matrix.

This is half of the simple confinement mechanism. At this point it is possible to obtain finite mass hadrons even though the parton masses diverge. However, since the cancellations are independent of the relative parton motion, we must study the residual interactions to see if they are confining. Since I am interested in the long-range interaction, I will study the fourier transform of the potential and compute $V(r) - V(0)$ so that the divergent constant in which we are no longer interested is canceled.

$$V_{singular}(r) - V_{singular}(0) \to \frac{g_\Lambda^2 C_F \Lambda^2}{4\pi^2 \mathcal{P}^+} \; \log\big(|x^-|\big) \; , \tag{19}$$

when $x_\perp = 0$ and $|x^-| \rightarrow \infty$; and

$$V_{singular}(r) - V_{singular}(0) \rightarrow \frac{g_\Lambda^2 C_F \Lambda^2}{2\pi^2 \mathcal{P}^+} \, \log\bigl(|x_\perp|\bigr) \;, \tag{20}$$

when $|x_\perp| \rightarrow \infty$ and $x^- = 0$. This potential is not rotationally symmetric, but it diverges logarithmically in all directions.

If the potential is not rotationally symmetric, how can rotational symmetry be restored? In light-front field theory rotations are dynamical. While it may be possible for rotational symmetry to be realized approximately in low-lying quark-antiquark states, exact symmetry requires additional explicit partons and even approximate rotational symmetry will require additional partons if we study highly excited states. We expect excited physical states in which a quark and antiquark are separated by a large distance to contain gluons. There is no reason to assume that the gluon content of these states is the same when the state is rotated, so rotational symmetry will be restored in highly excited states only if we allow additional partons. This complicates our attempt to derive a constituent picture, but we only need the constituent picture to work well for low-lying states. The intermediate range part of the potential is rotationally symmetric, and we may expect the ground state hadrons to be dominated by the valence configuration.

Isn't the confining potential supposed to be linear and not logarithmic? There is no conclusive evidence that the long-range potential is linear, and heavy quark phenomenology shows that a logarithmic potential can work quite well; however, lattice calculations provide strong evidence for a linear potential. However, low-lying states are not sensitive to the longest-range part of the interaction; and light quark-antiquark pairs prevent even exited states from being sensitive to the longest-range part of the interaction. In any case, I do not want to argue that these calculations show that the long-range potential in light-front QCD is logarithmic. Higher order corrections could produce powers of logarithms that add up to produce a linear potential.

The above argument seems to apply directly to QED at first sight. Is QED confining? There is a confining interaction between charged particles in the hamiltonian, but there is no strong interaction between charged particles and photons. To see if confinement survives in QED we should include the confining interaction in $H_0$, and then compute corrections in bound state perturbation theory. In QED the second order correction from photon exchange below the cutoff exactly cancels confinement. This implies that if confinement is included in $H_0$, higher order corrections are large. If the Coulomb interaction, which appears in $H$ also, is included in $H_0$, bound state perturbation theory appears to converge rapidly. On the other hand, in QCD gluons also experience a confining interaction. When second order bound state perturbation theory is used to study the effect of the exchange of confined gluons, it is seen that gluon exchange does not cancel the confining interaction in $H_0$; so this picture of confinement is at least self-consistent.

The important point is that $H$ contains a confining interaction that we are free to include in $H_0$, giving us some hope of finding a reasonable bound state perturbation theory for hadrons that resembles the bound state perturbation theory that has been

successfully applied to the study of atoms.

**Summary**

A constituent picture of hadrons may emerge in QCD if we use:

- hamiltonian light-front field theory
- a cutoff of order $\Lambda_{QCD}$ on energy changes that violates manifest covariance and gauge invariance
- a similarity renormalization group and coupling coherence

Bound states can be studied using bound state perturbation theory. The effective hamiltonian is computed as an expansion in the strong coupling, and then divided into $H = H_0 + V$, with $V$ treated perturbatively. $H_0$ must include all essential interactions. We have found that $H$ contains an order $\alpha$ logarithmically confining two-body interaction between all colored partons, and we have begun studies of bound states using this confining interaction, as discussed in the talk by Martina Brisudová.

**Acknowledgment**

I wish to thank the organizers for the opportunity to visit the beautiful Slovak Republic. This talk could be viewed as an introduction to Martina Brisudová's talk, which covered results for heavy mesons and I am indebted to Martina for many insights. I have profited from conversations with many people, particularly Ken Wilson, Stan Głazek, Brent Allen and Billy Jones. Finally I want to apologize to the many people whose work should have been referenced. This work was supported by the National Science Foundation under grant PHY-9409042.

**References**

1. K. G. Wilson, T. S. Walhout, A. Harindranath, W.-M. Zhang, R. J. Perry, St. D. Głazek: *Phys. Rev. D* **49** (1994) 6720;
2. R. J. Perry, in *Proceedings of Hadrons 94*, World Scientific, 1995 (hep-th/9407056);
3. K. G. Wilson and D. G. Robertson, in *Proceedings of the Fourth International Workshop on Light-Front Quantization and Non-Perturbative Dynamics*, World Scientific, 1995 (hep-th/9411007);
4. R. J. Perry, in *Proceedings of the Fourth International Workshop on Light-Front Quantization and Non-Perturbative Dynamics*, World Scientific, 1995 (hep-th/9411037);
5. M. Brisudová and R. Perry, *Initial studies of bound states in light-front QCD*, to appear in Phys. Rev. D, 1996 (hep-ph/9511443);
6. R. J. Perry, A. Harindranath, and K. G. Wilson: *Phys. Rev. Lett.* **65** (1990) 2959;
7. R. J. Perry: *Ann. Phys.* **232** (1994) 116;
8. St. D. Głazek and K.G. Wilson: *Phys. Rev. D* **48** (1993) 5863;
9. St. D. Głazek and K.G. Wilson: *Phys. Rev. D* **49** (1994) 4214;
10. R. J. Perry and K. G. Wilson: *Nucl. Phys. B* **403** (1993) 587;
11. W. M. Zhang and A. Harindranath: *Phys. Rev. D* **48** (1993) 4868, 4881, 4903;